Neutrinos and Cosmology: a Lifetime relationship

Pasquale D. Serpico

Center for Particle Astrophysics, Fermi National Accelerator Laboratory, Batavia, IL 60510-0500 USA

E-mail: serpico@fnal.gov

Abstract. We consider the example of neutrino decays to illustrate the profound relation between laboratory neutrino physics and cosmology. Two case studies are presented: In the first one, we show how the high precision cosmic microwave background spectral data collected by the FIRAS instrument on board of COBE, when combined with Lab data, have greatly changed bounds on the radiative neutrino lifetime. In the second case, we speculate on the consequence for neutrino physics of the cosmological detection of neutrino masses even as small as $\sim$0.06 eV, the lower limit guaranteed by neutrino oscillation experiments. We show that a detection at that level would improve by many orders of magnitude the existing limits on neutrino lifetime, and as a consequence on some models of neutrino secret interactions.

FERMILAB-CONF-08-144-A

1. Introduction

Historically, there has been a deep relation between neutrino physics and cosmology: the bounds on the number of generations from primordial nucleosynthesis, the neutrinos as dark matter candidates (and their effect on structure formation), or their masses as related to the origin of baryon asymmetry represent well known examples of this link. Here we take the example of neutrino decays to illustrate the interplay of the Lab and Cosmology in the neutrino sector with two case studies: In the first case, we show how the high precision cosmic microwave background spectral data collected by the FIRAS instrument on board of COBE, when combined with data from neutrino oscillation experiments and direct bounds on absolute masses, have greatly improved the bounds on the radiative neutrino lifetime. The relevant formalism is summarized in Sec. 2, while in Sec. 3 we present the bounds obtained. In the second case, presented in Sec. 4, we speculate on the consequence for neutrino physics of the cosmological detection of neutrino masses even as small as $\sim$0.06 eV, the lower limit guaranteed by neutrino oscillation experiments. We show that a detection at that level would improve by many orders of magnitude the existing limits on neutrino lifetime, and as a consequence on neutrino secret interactions with (quasi-)massless particles as in majoron models. In Sec. 5 we present our conclusions. For details on the two cases see respectively [1] and [2], on which this article is mostly based.

2. Radiative neutrino decays

Traditionally, constraints on neutrino radiative lifetime coming from cosmology were based on the diffuse Cosmic Infrared Background (CIB) and assumed strongly hierarchical masses in the eV range [3, 4]. These constraints are now outdated and strictly speaking inapplicable. The neutrino mass splittings squared provided by oscillation experiments and present upper bounds on the neutrino mass scale constrain neutrino mass differences to fall in the microwave energy range ($E \sim 10^{-3}$ eV), in most of the allowed parameter space. This implies that current bounds can take advantage of the high precision cosmic microwave background (CMB) data collected by the FIRAS instrument on board of COBE, which tested the blackbody nature of the spectrum at better than 1 part in 10^4 [5, 6].

Let us denote by ν_i the neutrino fields respectively of masses m_i, where $i = 1, 2, 3$. The radiative decay $\nu_i \to \nu_j + \gamma$ can be thought of as arising from an *effective* interaction Lagrangian of the form

$$\mathcal{L}_{\rm int} = \frac{1}{2}\bar{\nu}^i \sigma_{\alpha\beta}(\mu_{ij} + \epsilon_{ij}\gamma_5)\nu^j F^{\alpha\beta} + \text{ h.c.} \tag{1}$$

where $F^{\alpha\beta}$ is the electromagnetic field tensor, $\sigma_{\alpha\beta} = [\gamma_\alpha, \gamma_\beta]$ where γ_μ are the Dirac-matrices and $[.,.]$ is the commutator and μ_{ij} and ϵ_{ij} are the magnetic and electric transition moments usually expressed in units of the Bohr magneton μ_B. The convention to sum over repeated indices is used. In general, μ_{ij} and ϵ_{ij} are functions of the transferred momentum squared q^2, so that constraints obtained at a different q^2 are independent. The radiative decay rate for a transition $i \to j$ is written

$$\Gamma_{ij}^{\gamma} = \frac{|\mu_{ij}|^2 + |\epsilon_{ij}|^2}{8\pi}\left(\frac{m_i^2 - m_j^2}{m_i}\right)^3 \equiv \frac{\kappa_{ij}^2}{8\pi}\left(\frac{m_i^2 - m_j^2}{m_i}\right)^3 .$$

In the following, we shall assume that the radiative decay rate is very low compared with the expansion rate of the universe; neither the cosmological evolution or the primordial neutrino spectrum is affected by the additional coupling we are going to introduce. A posteriori, this is known to be an excellent approximation. For the same reason, we shall also neglect "multiple decays". We shall take our input data for neutrino mass eigenstate densities from the calculation performed in [7] without any extra parameter, as non-vanishing chemical potentials. With present data, the latter are anyway constrained to be well below $\mathcal{O}(1)$ [8], so dropping this assumption would not change much our conclusions.

From simple kinematical considerations it follows that in a decay $\nu_i \to \nu_j + \gamma$ from a state of mass m_i into one of mass $m_j < m_i$, the photon in the rest frame of the decaying neutrinos is thus monochromatic (two-body decay), with an energy

$$\varepsilon_{ij} = \frac{m_i^2 - m_j^2}{2\, m_i} . \tag{2}$$

At present, the neutrino mass spectrum is constrained by the well-known values of the two squared mass splittings for the atmospheric (Δm_H^2) and the solar (Δm_L^2) neutrino problems. We take their best-fit values and 2σ ranges from [9], $\Delta m_L^2 = 7.92\,(1 \pm 0.09) \times 10^{-5}$ eV^2, $\Delta m_H^2 = 2.6\,(1^{+0.14}_{-0.15}) \times 10^{-3}$ eV^2 . The remaining unknowns in the neutrino spectrum are the absolute mass scale (equivalently, the mass of the lightest eigenstate m_1) and the mass hierarchy. Namely, in normal hierarchy (NH) the mass pattern would be $\{m_1\,, m_2 = \sqrt{m_1^2 + \Delta m_L^2}\,, m_3 = \sqrt{m_1^2 + \Delta m_L^2 + \Delta m_H^2}\,\}$ while in inverted hierarchy (IH) one would have $\{m_1, m_2 = \sqrt{m_1^2 + \Delta m_H^2}\,, m_3 = \sqrt{m_1^2 + \Delta m_L^2 + \Delta m_H^2}\}$.

In the limiting case of normal hierarchy and $m_1 = 0$, the lightest neutrino for which a decay is possible has a mass $m_2 \simeq 9 \times 10^{-3}$ eV and is non-relativistic for most of the universe lifetime, namely in the redshift range $z \lesssim 50$. We can thus safely work in the approximation of all neutrinos decaying effectively at rest. In this limit, we can also neglect the momentum distribution of the neutrino spectra. We shall vary the mass scale in $0 \lesssim m_1 \lesssim 2\,$eV as allowed by the Mainz experiment on the ^{3}H beta decay endpoint [10].

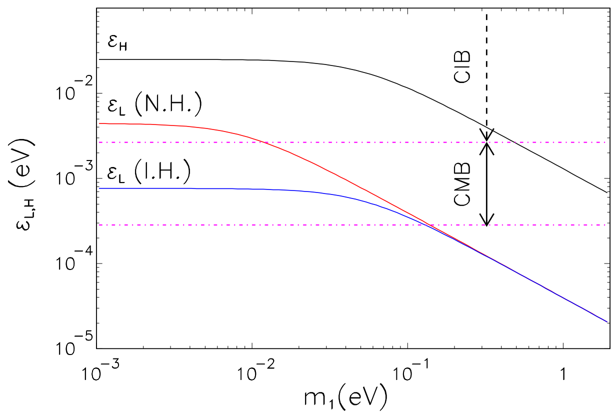

Figure 1. Unredshifted photon energy ε from decaying neutrinos [Eq. (2)] as a function of the lightest neutrino mass eigenstate m_1, for the two neutrino mass splittings (L,H) in normal and inverted hierarchy (see text for details). The horizontal band represents the energy range of the CMB spectrum measured by FIRAS [5]. The CIB energy range is also shown.

Let F_E be the present energy flux of photons with present energy E produced in the decay. The differential energy flux φ_E (energy flux F_E per unit energy and solid angle) is related to the differential number flux φ_n (the particle flux F_n per unit energy and solid angle) at present by

$$\varphi_E \equiv \frac{\mathrm{d}^2 F_E}{\mathrm{d}E\,\mathrm{d}\Omega} = E\,\frac{\mathrm{d}^2 F_n}{\mathrm{d}E\,\mathrm{d}\Omega} \equiv E\,\varphi_n, \tag{3}$$

and it can be shown that, if the lifetime τ_i of the neutrino of mass m_i is much greater than the universe lifetime it holds [11]

$$\varphi_E = \frac{\Gamma^\gamma_{32}}{4\pi}\frac{n_3}{H(z_{32})} + \frac{\Gamma^\gamma_{31}}{4\pi}\frac{n_3}{H(z_{31})} + \frac{\Gamma^\gamma_{21}}{4\pi}\frac{n_2}{H(z_{21})}\,, \tag{4}$$

where $n_i \simeq 113$ cm^{-3} is the present number density of the $i-$th neutrino *in absence of decay*, the Hubble function is (assuming, for simplicity, a flat cosmology) $H(z) = H_0\sqrt{\Omega_M(1+z)^3+\Omega_\Lambda}$, $H_0 \simeq 73$ km s^{-1} Mpc^{-1} being the present Hubble expansion rate, and $\Omega_M \simeq 0.26$ and $\Omega_\Lambda \simeq 0.74$ respectively the matter and the cosmological constant energy density relative to the critical one. The dependence on energy enters implicitly via the quantities $0 \le z_{ij} = \varepsilon_{ij}/E - 1$.

In practice, to a very good approximation one can write a general equation of the kind

$$\varphi_E = \frac{\Gamma^\gamma_H}{4\pi}\frac{n_H}{H(z_H)} + \frac{\Gamma^\gamma_L}{4\pi}\frac{n_L}{H(z_L)}\,, \tag{5}$$

where the meaning of the factors however depends on the hierarchy. In NH, in the first two terms of the sum in Eq. (4) it holds $z_{32} \simeq z_{31} \equiv z_H$, and one can identify $z_L = z_{21}$, $\Gamma^\gamma_L = \Gamma^\gamma_{21}$, $\Gamma^\gamma_H = \Gamma^\gamma_{31} + \Gamma^\gamma_{32}$. In IH, it is the last two terms of the sum in Eq. (4) which have $z_{31} \simeq z_{21} \equiv z_H$, and using $n_2 \simeq n_3$ one can identify $z_L = z_{32}$, $\Gamma^\gamma_L = \Gamma^\gamma_{32}$, $\Gamma^\gamma_H \equiv \Gamma^\gamma_{31} + \Gamma^\gamma_{21}$.

In Fig. 1 we represent the unredshifted photon energy ε_{ij} from decaying neutrinos [Eq. (2)] as a function of the lightest neutrino mass eigenstate m_1 in the case of normal and inverted hierarchy, where the meaning of $\varepsilon_{L,H}$ is clear from the previous discussion. We also indicate by an horizontal band the energy range of the CMB spectrum (2.84×10^{-4} eV $\le E \le 2.65\times10^{-3}$ eV) measured by FIRAS [5]. We also show the CIB range in the energy band above the FIRAS range up to (conventionally) 0.15 eV [12]. For $m_1 \lesssim 0.5$ eV, ε_H falls in the CIB range.

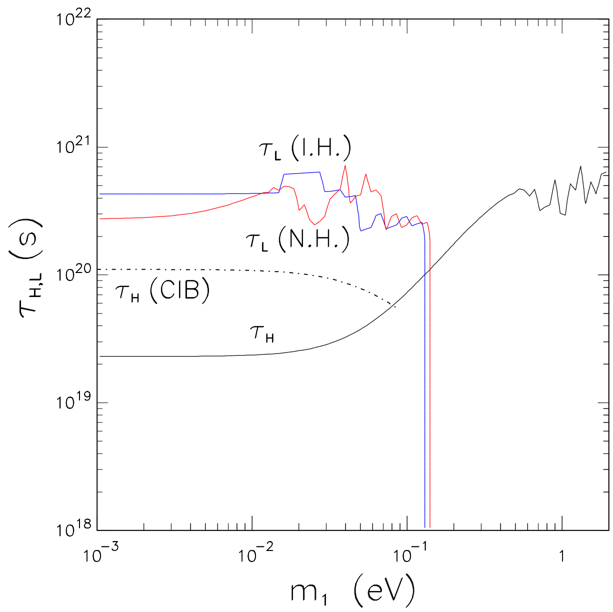

Figure 2. Bounds on τ_H and τ_L vs. m_1, for the two cases of NH and IH. The regions below the solid curves are excluded at 95 % C.L. The NH and IH curves for τ_H coincide, although the definition of τ_H is different (see text). The dot-dashed line is the limit obtained from the CIB.

For photons emitted at $z = 0$ in the FIRAS range, the effect of radiative decays is most prominent and results in a feature on the CMB spectrum. Actually even if photons are emitted at higher energy the effect is still strong, since photons emitted at a redshift of a few enter the FIRAS spectrum because of cosmological redshift; it is easy to check that one has thus some sensitivity to κ_H in the whole range for m_1. On the other hand, for $m_1 \gtrsim 0.14$ eV the photons corresponding to the smaller splitting are falling in the radio band, below the frequency range probed by COBE, where measurements are more uncertain and thus one has no sensitivity to κ_L and the corresponding bound disappears.

3. The radiative decay bound

To constrain the neutrino electromagnetic decay we use the COBE/FIRAS data for the experimentally measured CMB spectrum, corrected for foregrounds [5, 6]. The $N = 43$ data points $\Phi_i^{\rm exp}$ at different energies E_i are obtained by summing the best-fit blackbody spectrum to the residuals reported in Ref. [5]. The experimental errors σ_i and the correlation indices ρ_{ij} between different energies are also available. In the presence of neutrino decay, the original radiance of the "theoretical blackbody" at temperature T

$$\Phi^0(E,T) = \frac{E^3}{4\pi^3}\left[\exp(E/T) - 1\right]^{-1} \tag{6}$$

would gain an additional term so that the intensity becomes

$$\Phi^0(E,T) \quad \rightarrow \quad \Phi(E,T,\kappa^2_{L,H},m_1) = \Phi^0(E,T) + \varphi_E(\kappa^2_{L,H},m_1).$$

The bound are derived via a χ^2 technique, taking into account the covariance matrix of the data points and leaving the parameter T entering initially in $\Phi^0(E,T)$ free. Indeed, it needs not to be fixed at the standard value $T_0 = 2.725 \pm 0.002$ K [6], which is the best fit of the "distorted" spectrum eventually observed now, but is left to be determined in the minimization procedure.

In Fig. 2 we report the exclusion plot in the plane $\tau_{L,H} \equiv (\Gamma^{\gamma}_{H,L})^{-1}$ vs. m_1, where the regions below the solid curves are excluded at 95 % C.L. For small values of m_1 the most stringent limit is $\tau_L \gtrsim 4 \times 10^{20}$ s in IH (slightly better than in NH case), while the bound on τ_H is about an order of magnitude smaller, say $\tau_H \gtrsim 2 \times 10^{19}$ s, since for low m_1 only photons produced by H decays at a redshift of few are in FIRAS range. On the contrary, for $m_1 \gtrsim 0.14$ eV, the bound on τ_L disappears, while the bound on τ_H becomes more stringent, being $\tau_H \gtrsim 5 \times 10^{20}$ s. The "fuzzy" behaviour of the bounds is due to the sharp edge of the photon spectrum at $E = \varepsilon_{H,L}$: when the photon energy embeds a new FIRAS bin, the χ^2 function has a sharp discontinuity. If one translates the bounds of Fig. 2 into bounds on $\kappa_{L,H}$, the factor $(\delta m^2_{ij}/m_i)^3$ plays a significant role. For the NH case, $\kappa_L \lesssim 3 \times 10^{-8}\,\mu_B$, while in the IH case, $\kappa_L \lesssim 3 \times 10^{-7}\,\mu_B$. In agreement with our previous considerations, the bound on κ_L disappears for $m_1 \gtrsim 0.14\,$eV. On the contrary, the bound for κ_H is always present, and it corresponds to $\kappa_H \lesssim 8 \times 10^{-9}\,\mu_B$ apart for the degenerate region, where it degrades down to $10^{-7}\,\mu_B$ or even more.

It is worth noting that an improved bound on τ_H for low m_1 can be obtained from observations of the CIB, which differently from the CMB does not origin in the early universe, being rather the relic emission of all the galaxies at wavelengths larger than a few microns. Recently, a new estimate of the CIB flux has been established using the Spitzer Observatory data [12]. Using this determination one can obtain a rough bound on τ_H simply requiring that the total energy flux of the photons coming from ν decay does not exceed the CIB flux:

$$\int_{E_{\min}}^{\varepsilon_H} \varphi_E\, dE < \Phi_{\rm CIB} \sim 24\,{\rm nW\,m^{-2}\,sr^{-1}} \quad , \tag{7}$$

where we consider as lower limit of the CIB range the upper value of the FIRAS range, i.e. $E_{\min} = 2.65 \times 10^{-3}$ eV. The bounds on τ_H from Eq. (7) is shown in Fig. 2 by the dot-dashed line. Although this bound is stronger than those obtained by the FIRAS data in the same range of m_1, we emphasize that it should be considered only as indicative, due to the larger uncertainties in the CIB normalization and spectral shape. Interestingly, if it turns out that $m_1 \lesssim 0.1\,$eV, an improvement on the bound on $\tau_{\rm H}$ will clearly take advantage of a better measurement of the CIB flux and a more detailed knowledge of the astrophysical sources contributing to it.

4. Invisible decays and secret neutrino interactions

Recent years have seen an impressive improvement on the cosmological constraints to the sum of neutrino masses $\Sigma = \sum m_i$ (for reviews see [13, 14]), with current limit from CMB only at the level of 1.3 eV [15], while combinations of different datasets produce bounds below 1 eV, with the most aggressive bounds (but also the most fragile ones with respect to unaccounted systematics) already at $\Sigma \lesssim 0.2$ eV, 95% C.L. [16, 17]. Several forecast analyses suggest that cosmological probes will reach in the future an incredible sensitivity to the effects of even a tiny mass of the cosmic background neutrinos. In particular, cosmic microwave background (CMB) lensing extraction may be sensitive to $\Sigma \simeq 0.035\,$eV [18]; CMB plus weak galaxy lensing with tomography may also push the sensitivity to Σ below the level of $\sim 0.05\,$eV [19, 20], with an error as low as $\sim 0.013\,$eV [19]. Also galaxy cluster surveys may probe Σ down to $\sim 0.03\,$eV [21], and a sensitivity down to $0.05 \pm 0.015\,$eV may be reached combining CMB with the data from the Square Kilometre Array survey of large scale structures [22]. These forecasts show that cosmology has a potential sensitivity to neutrino masses well below the 0.1 eV level. Of course, the ultimate level of the systematics to beat has yet to be reliably established. On the other hand, the synergy between different strategies and probes may help to identify the systematics, and also to improve over the above-mentioned figures of merit.

The interest of these expectations relies on the fact that neutrino oscillation data imply $\Sigma \gtrsim 0.06\,$eV, where the minimum $\Sigma \simeq 0.061 \pm 0.004\,$eV is attained for a normal hierarchy (NH; values quoted at 2 σ, see [17]). For the case of an inverted mass hierarchy (IH), the oscillation data imply $\Sigma \simeq 0.1\,\mathrm{eV}$. The following arguments assume that neutrinos have a hierarchical spectrum of either inverted or normal sign, as favored by many theoretical models, including the simplest seesaw ones. We argue that, if a positive cosmological mass detection is achieved as expected, one will be able to put a remarkably strong constraint on the neutrino lifetime. Note that previous attention has been paid to the cosmological signatures of decaying neutrinos [23]. Yet the mass range explored in those papers is very large compared to present bounds, and the main signature considered was the impact on the integrated Sachs-Wolfe effect on the CMB. In our considerations, the bound comes from the impact that massive neutrinos have in the background evolution of the universe, in a range of masses where they are relativistic well after the CMB decoupling.

Bounds on neutrino lifetimes are usually quoted in terms of the rest-frame lifetime to mass ratio τ/m. Given a measurement in the time interval t using neutrinos with Lab energy E, the naive bound which one can put is $\tau/m \gtrsim t/E$. Using then the longest timescale available, the universe lifetime $t_0 \simeq H_0^{-1}$ (where H_0 is the Hubble constant), and the lowest energy neutrinos, the ones of the cosmic background which are at least partially non-relativistic, a bound of the order of ($m_{50} \equiv m/50\,\mathrm{meV}$)

$$\frac{\tau}{m} \gtrsim \frac{1}{m\,H_0} \simeq 10^{19}\, m_{50}^{-1}\,\mathrm{s/eV}\,, \tag{8}$$

is the strongest constraint attainable in principle. This is to be compared with the strongest direct bound available at present given by the observation of solar MeV neutrinos, of the order of $\sim 10^{-4}\,\mathrm{s/eV}$ [24]. A bound of the order of $\tau/m \gtrsim 4 \times 10^{11} m_{50}^2\,\mathrm{s/eV}$, has been claimed to follow already from the requirement that the neutrinos are free-streaming at the time of the photon decoupling, as deduced by precise measurements of the CMB acoustic peaks [25]. Yet, the robustness of this conclusion has been questioned in [26]. We shall see that the proposed bound based on cosmological neutrino mass detection would be much closer to the maximal theoretical bound of Eq. (8), thus superseding by several orders of magnitude the previous ones. More importantly, it is not based on a model for secret neutrino interactions, but on the "observation" of neutrino survival, and it applies whatever the final state light particles are.

Let us define the fraction of matter energy density in neutrinos as $f_\nu \equiv \rho_\nu/(\rho_{\mathrm{m}} + \rho_\nu)$, where ρ_{m} is the average cold dark matter (CDM) plus baryon density, and ρ_ν is the neutrino energy density. Keeping e.g. $\rho_{\mathrm{m}} + \rho_\nu$ constant, a non-vanishing f_ν today would change the redshift of matter radiation equality z_{eq} with respect to the massless neutrino case, as well as attenuate the growth of structures. The combined effect of the shift in the time of equality and of the reduced CDM fluctuation growth during matter domination produces an attenuation of perturbations for modes $k > k_{\mathrm{nr}}$, where k_{nr} is the minimum of the comoving free-streaming wavenumber attained when neutrinos turn non-relativistic, and given by [13]

$$k_{\mathrm{nr}} \simeq 1.5 \times 10^{-3} m_{50}^{1/2}\,\mathrm{Mpc}^{-1}. \tag{9}$$

An instantaneous decay of the massive neutrinos at a redshift z_d in the matter era can be thought as replacing the neutrino fluid with one having the same energy content at z_d, but whose energy density scales from that moment on as $(1+z)^4$, since the daughter particles are relativistic. Let us estimate how large a value of z_d, or equivalently of the proper time $t_d (= \tau$ if the neutrino is non-relativistic), can be probed cosmologically. Quickly after the neutrino decay one has formally $f_\nu \to 0$, provided that $t_d \ll t_0 \simeq H_0^{-1}$; from that moment on, the cosmological effects of the decaying neutrino scenario are analogous to the ones of a massless neutrino universe. The condition $t_d \ll t_0$ is required by the fact that when $t_d \to t_0$, the radiation content of the relativistic daughters of the massive neutrino has no time to decline to zero with respect to the matter density. This condition is necessary to change appreciably the energy budget of the universe, thus affecting the predicted growth of the structures and the time of equality with respect to a massive neutrino scenario. Clearly, for a given sensitivity to the effect of neutrino masses there is a maximum value $t_d^{\rm max}$ which would result in a detectable change of cosmological observables. A precise estimate of this parameter would imply a detailed forecast analysis, which goes beyond the purpose of this paper. Yet, a simple argument shows that, relying on the existing forecasts, a conservative lower limit is $t_d^{\rm max} \gtrsim t_{\rm nr}$, where $t_{\rm nr}$ is the epoch at which the heavier neutrinos become non-relativistic, whose redshift is defined by $m = 3\,T_{\nu,0}(1+z_{\rm nr})$, $T_{\nu,0}$ being the present temperature of the neutrino gas. Indeed, when the decay epoch satisfies $t_{\rm d} \lesssim t_{\rm nr}$, the energy content of the products is the same of a relativistic neutrino fluid, and it redshifts the same way. So, all physical effects of this scenario are basically the same of the case where neutrinos are massless. In Fig. 3, from top to bottom as seen from the left side of the plot, we show $f_\nu(z)$ for the following cases: (i) a massive neutrino cosmology, where we assume an IH neutrino mass pattern and the lightest neutrino is massless; (ii) as in (i), but for NH; (iii) a decaying neutrino cosmology, where massive neutrinos have IH; (iv) as in (iii), but for NH; (v) a massless neutrino cosmology. For the decaying cases, we assume that all massive neutrinos decay at $t_d = t_{\rm nr}$, where $t_{\rm nr}$ is the time of non-relativistic transition of the heaviest neutrino state ($m \simeq 0.05\,{\rm eV}$) . For simplicity we have assumed a matter-dominated cosmology with the matter density parameter $\Omega_{\rm m} = 0.24$ and the reduced Hubble constant $h = 0.73$ [27].

Clearly, the cases (iii), (iv), and (v) are very similar (exactly degenerate if $t_d \ll t_{\rm nr}$) and, as long as $t_{\rm d} \lesssim t_{\rm nr}$, if the massless neutrino case can be disproved, the decaying neutrino bound immediately follows. The improvement in the bound on the neutrino lifetime is tremendous. In particular, neutrinos turn non-relativistic at $z_{\rm nr} \simeq m/3\,T_{\nu,0} \simeq 100\,m_{50}$, i.e. when the universe has about $(100\,m_{50})^{-3/2} \sim 10^{-3} m_{50}^{-3/2}$ of its present age, and the bound is about $10^{-3} m_{50}^{-3/2}$ of the maximum attainable limit reported in Eq. (8),

$$\frac{\tau}{m} \gtrsim 10^{16}\, m_{50}^{-5/2}\ {\rm s/eV}\,. \tag{10}$$

Obviously, the previous argument does not exclude that an accurate forecast analysis may reveal a sensitivity to a somewhat larger $t_d^{\rm max}$. Note also that we do not require that cosmological data need to distinguish between NH and IH: if future observations will suggest e.g. $\Sigma = 0.08\,{\rm eV}$ with a 1σ error of 0.02 eV, the two neutrino mass patterns would be both consistent within $1\,\sigma$ with the best fit, yet a complete decay of neutrinos into relativistic particles with lifetime lower than the value reported in Eq. (10) could be excluded at $4\,\sigma$. Of course, for a given cosmological sensitivity, the significance of the above bound increases if the inverted hierarchy is realized in nature: in that case $\Sigma \simeq 0.1\,{\rm eV}$ holds, and the cosmological effects of neutrino masses are larger. Note that accelerator neutrino experiments, magnetized detectors of atmospheric neutrinos, direct mass searches, and the serendipitous observation of neutrinos from a galactic supernova may all be used to determine the mass hierarchy. It is thus possible that by the time cosmology will be sensitive to $\Sigma \lesssim 0.1\,{\rm eV}$, the hierarchy information may be available independently.

To appreciate how strong the bound of Eq. (10) would be, let us consider a model of a "secret" neutrino interaction with a (quasi-)massless majoron field ϕ of the kind $\mathcal{L} = g\,\bar\nu_i\nu_j\phi+$ h.c, i,j labeling different mass eigenstates. The total decay rate for a hierarchical neutrino mass pattern and summing over neutrino and antineutrino final state channels is [24, 25]

$$\Gamma_{\rm d} = t_d^{-1} = \frac{g^2}{16\pi} m\,. \tag{11}$$

This holds in the neutrino rest frame, but in our case this is also the Lab decay width, give or take a factor $\mathcal{O}(1)$, since the neutrino is just turning non-relativistic. The constraint of Eq. (10) leads to the stringent bound

$$g \lesssim 4\times 10^{-14}\, m_{50}^{1/4}. \tag{12}$$

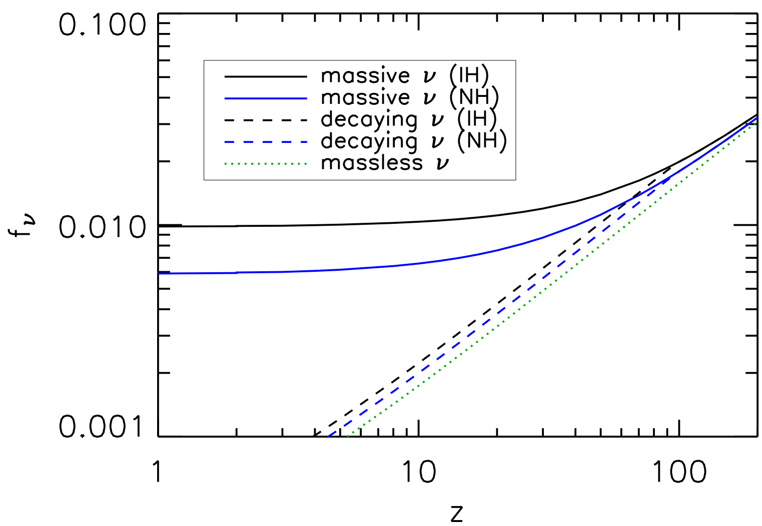

Figure 3. The function $f_\nu(z)$ for the relevant cosmological cases considered in the text.

This has to be compared with traditional bounds found in the literature in the range $g \lesssim 10^{-4} \div 10^{-5}$ (see e.g. [28]). Even the extremely stringent bound reported in [25] is more than two orders of magnitude weaker. Note that the tiny couplings which may induce the decay are not sufficient to thermalize extra degrees of freedom in the early universe. So, this model does not predict departure from the standard expectation for the effective number of neutrinos $N_{\rm eff}$ which can be consistently fixed in deriving the bound.

5. Discussion and Conclusions

In this paper, we have revisited the bound on the neutrino radiative lifetime coming from cosmology, deriving updated constraints from the high precision CMB spectrum data collected by the FIRAS instrument on board of COBE, within the presently allowed range of mass parameters suggested by neutrino oscillation physics and tritium endpoint experiments. The updated bounds are reported in Fig. 2. The improvement over pre-existing bounds depends both on narrowing the uncertainty on the neutrino mass spectrum and on the observational test of the blackbody nature of the spectrum at better than 1 part in 10^4 [5, 6]. This provides and explicit example of the continuing interplay of Lab neutrino experiments and cosmology.

Looking at the future, it is very interesting that current forecast analyses suggest that future cosmological surveys may attain the sensitivity to detect the effects of a sum of neutrino masses as small as $\sim$0.06 eV, the lower limit predicted by oscillation data. Provided that the systematics can be controlled to that level, we have discussed in this paper how such a detection would have profound consequences for the particle physics of the neutrino sector, besides providing a way to measure the absolute neutrino mass scale. In particular, when taking into account the expectations from the Lab, excluding the $\Sigma = 0$ case would improve by many orders of magnitude the existing limits on neutrino lifetime, and as a consequence on neutrino secret interactions with (quasi-)massless particles as in majoron models. Strictly speaking these bounds apply to the heaviest (or the two heaviest, in IH) mass eigenstate, but naturaleness and phase-space considerations suggest that the lifetime of the lightest state(s) is longer, and its coupling with a majoron field weaker, than for the heavier one(s). It also applies to any possible invisible decay channel, provided that the total mass of the final state particles is much smaller than Σ. In particular, this bound applies to 3-ν final state decays $\nu_i \to \bar{\nu}_j \nu_j \nu_k$, as well as to decays $\nu_i \to \nu_j + \phi$ in majoron-like models. We think that this idea summarized here provides another beautiful example of interplay between particle physics and cosmological arguments and motivates further the efforts to fully exploit the potential of future cosmological surveys.

Acknowledgments

I would like to thank the organizers for the kind invitation to celebrate the 75th birthdays of such two prominent members in the field of neutrino and low background physics, Frank Avignone and Ettore Fiorini. I would also like to dedicate this paper to the memory of S. Peter Rosen, an inspirational figure for all neutrino physicists. It is a pleasure to acknowledge A. Mirizzi and D. Montanino for collaborating on part of the topics summarized here. Support by the US Department of Energy and by NASA grant NAG5-10842 is acknowledged. Fermilab is operated by Fermi Research Alliance, LLC under Contract No. DE-AC02-07CH11359 with the United States Department of Energy.

References

[1] A. Mirizzi, D. Montanino and P. D. Serpico, Phys. Rev. D **76**, 053007 (2007).
[2] P. D. Serpico, Phys. Rev. Lett. **98**, 171301 (2007).
[3] M.T. Ressell and M.S. Turner, Comm. Astrophys. **14**, 323 (1990).
[4] S. D. Biller *et al.*, Phys. Rev. Lett. **80**, 2992 (1998).
[5] D. J. Fixsen *et al.*, Astrophys. J. **473**, 576 (1996).
[6] J. C. Mather *et al.* Astrophys. J. **512**, 511 (1999).
[7] G. Mangano *et al.*, Nucl. Phys. B **729**, 221 (2005).
[8] P. D. Serpico and G. G. Raffelt, Phys. Rev. D **71**, 127301 (2005).
[9] G. L. Fogli, E. Lisi, A. Marrone and A. Palazzo, Prog. Part. Nucl. Phys. **57**, 742 (2006).
[10] C. Kraus *et al.*, Eur. Phys. J. C **40**, 447 (2005).
[11] E. Massò and R. Toldrà, Phys. Rev. D **60**, 083503 (1999).
[12] H. Dole *et al.*, astro-ph/0603208.
[13] J. Lesgourgues and S. Pastor, Phys. Rept. **429**, 307 (2006).
[14] S. Hannestad, Ann. Rev. Nucl. Part. Sci. **56**, 137 (2006).
[15] J. Dunkley *et al.* [WMAP Collaboration], arXiv:0803.0586 [astro-ph].
[16] U. Seljak, A. Slosar and P. McDonald, JCAP **0610**, 014 (2006).
[17] G. L. Fogli *et al.*, Phys. Rev. D **75**, 053001 (2007).
[18] J. Lesgourgues, L. Perotto, S. Pastor and M. Piat, Phys. Rev. D **73**, 045021 (2006).
[19] Y. S. Song and L. Knox, Phys. Rev. D **70**, 063510 (2004).
[20] S. Hannestad, H. Tu and Y. Y. Y. Wong, JCAP **0606**, 025 (2006).
[21] S. Wang, Z. Haiman, W. Hu, J. Khoury and M. May, Phys. Rev. Lett. **95**, 011302 (2005).
[22] F. B. Abdalla and S. Rawlings, astro-ph/0702314.
[23] R. E. Lopez, S. Dodelson, R. J. Scherrer and M. S. Turner, Phys. Rev. Lett. **81**, 3075 (1998); S. Hannestad, Phys. Rev. D **59**, 125020 (1999); M. Kaplinghat, R. E. Lopez, S. Dodelson and R. J. Scherrer, Phys. Rev. D **60**, 123508 (1999); M. J. White, G. Gelmini and J. Silk, Phys. Rev. D **51**, 2669 (1995).
[24] J. F. Beacom and N. F. Bell, Phys. Rev. D **65**, 113009 (2002).
[25] S. Hannestad and G. Raffelt, Phys. Rev. D **72**, 103514 (2005); See also A. Basboll, O. E. Bjaelde, S. Hannestad and G. G. Raffelt, arXiv:0806.1735 [astro-ph].
[26] N. F. Bell, E. Pierpaoli and K. Sigurdson, Phys. Rev. D **73**, 063523 (2006).
[27] W. M. Yao *et al.* [Particle Data Group], J. Phys. G **33**, 1 (2006).
[28] Y. Farzan, Phys. Rev. D **67**, 073015 (2003).